New inflation in supergravity with a chaotic initial condition

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Abstract

We propose a self-consistent scenario of new inflation in supergravity. Chaotic inflation first takes place around the Planck scale, which solves the longevity problem, namely, why the universe can live much beyond the Planck time, and also gives an adequate initial condition for new inflation. Then, new inflation lasts long to generate primordial fluctuations for the large scale structure, which generally has a tilted spectrum with the spectral index $n_s < 1$. The successive decay of the inflaton leads to the reheating temperature low enough to avoid the overproduction of gravitinos in a wide range of the gravitino mass.

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I. INTRODUCTION

Inflation is the most natural extension of the standard big-bang cosmology because it not only solves the flatness and the horizon problems but also provides a generation mechanism of density fluctuations [1–3]. Since inflation typically takes place at an energy scale much higher than the electroweak scale, most favorably around the Planck scale in fact, we cannot but seriously confront with the hierarchy problem in particle physics. To this end, it is inevitable to consider inflation models in the supersymmetric theory, in particular, its local version, supergravity (SUGRA) [4,5].

Among various mechanisms of inflation proposed so far, chaotic inflation [6] is the most attractive in that it starts around the Planck scale so that it does not suffer from the longevity problem, that is, why the universe lives beyond the Planck time. However, it has been recognized to be difficult to realize chaotic inflation in SUGRA because scalar potentials in minimal SUGRA have an exponential factor of the form $\exp(|\phi|^2/M_G^2)$ which forbids any scalar field ϕ to have a value much larger than the reduced Planck scale $M_G \simeq 2.4\times 10^{18}$ GeV and hampers chaotic inflation. Although several supergravity chaotic inflation models have been proposed so far using functional degrees of freedom in SUGRA as a non-renormalizable theory [7,8], these models employ rather specific Kähler potentials and require fine tuning since there are no symmetry reasons for having the proposed forms. Recently, however, Kawasaki, Yanagida, and one of the authors (M.Y.) [9] proposed a natural model of chaotic inflation in SUGRA using the Nambu-Goldstone-like shift symmetry. But, in this model, the reheat temperature low enough to avoid overproduction of gravitinos is realized taking a rather small coupling constant ($\sim 10^{-5}$) though it is natural in 't Hooft's sense [10].

As a model of inflation that predicts low reheating temperature straightforwardly [11–13], new inflation is more attractive than chaotic inflation because it occurs at a lower energy scale. Furthermore, new inflation can easily generate density fluctuations [14] with a tilted spectrum, which may naturally explain the recent observation of anisotropies of the cosmic microwave background radiation (CMB) by the BOOMERANG experiment [15] and the MAXIMA experiment [16]. However, it suffers from a severe problem about the initial value of the inflaton [3] in addition to the longevity problem mentioned above. In order to realize successful inflation, the initial value of the inflaton must be fine-tuned near the local maximum of the potential over the horizon scale. For this problem Asaka, Kawasaki, and one of the authors (M.Y.) [17] proposed a solution by considering the gravitationally suppressed interactions with particles in the thermal bath. But the longevity problem remains. Izawa, Kawasaki, and Yanagida [18], on the other hand, considered another inflation (called pre-inflation) which takes place before new inflation and drives the scalar field responsible for new inflation dynamically toward the local maximum of its potential. If the pre-inflation is chaotic inflation, the longevity problem is solved, too.

Thus we are naturally motivated to a model of successive inflation, namely, chaotic inflation followed by new inflation. In fact, such a double inflation model has already been proposed in a different context [19], but in this paper, we propose a simple and self-consistent model of successive inflation in the framework of SUGRA in which the two inflatons belong to the same supermultiplet. That is, the inflaton for chaotic inflation is the imaginary part of a complex scalar field while its real part drives new inflation. In fact, our model is a triple inflation model where chaotic inflation first takes place followed by a mini inflation driven by a false vacuum energy and then new inflation occurs.

II. MODEL

We introduce the inflaton chiral superfield $\Phi(x, \theta)$ with the Nambu-Goldstone-like shift symmetry as considered in ref. [9]. That is, the Kähler potential $K(\Phi, \Phi^*)$ is invariant under the following shift of Φ,

$$\Phi \to \Phi + i\; CM_G, \tag{1}$$

where C is a dimensionless real constant. Then, the Kähler potential $K(\Phi, \Phi^*)$ becomes a function of $\Phi + \Phi^*$, which allows the imaginary part of the scalar component of the superfield Φ to take a value larger than M_G without costing exponentially large potential energy density. Hereafter we set the reduced Planck scale to be unity.

As far as the Nambu-Goldstone-like shift symmetry is exact, the inflaton cannot have a potential. We, therefore, propose the following superpotential which breaks the shift symmetry with a small dimensionless coupling parameter g',

$$W = v^2 X - g' X \Phi^2 = v^2 X (1 - g\Phi^2), \tag{2}$$

where we have introduced another chiral superfield $X(x, \theta)$. Here v is a scale generated dynamically, and $g \equiv g' v^{-2}$. This superpotential W is invariant under $U(1)_R \times Z_2$ symmetry. Under $U(1)_R$ symmetry,

$$X(\theta) \to e^{-2i\alpha} X(\theta e^{i\alpha}), \tag{3}$$

$$\Phi(\theta) \to \Phi(\theta e^{i\alpha}). \tag{4}$$

Also, X is even and Φ is odd under Z_2 symmetry. Although this superpotential W is not invariant under the Nambu-Goldstone-like shift symmetry, it is natural in the sense of 't Hooft [10]. For the symmetry is recovered if the small parameter g' is set to be zero. Since the correction to the Kähler potential from this breaking term is negligible if $g \ll 1$, we consider the Kähler potential which is invariant under the Nambu-Goldstone-like shift symmetry and $U(1)_R \times Z_2$ symmetry,

$$K(\Phi, \Phi^*, X, X^*) = K[(\Phi + \Phi^*)^2, XX^*]. \tag{5}$$

Below we take a minimal Kähler potential, for which scalar kinetic terms are canonical,

$$K = \frac{1}{2}(\Phi + \Phi^*)^2 + XX^*, \tag{6}$$

and extend it later to incorporate higher-order effects.

III. DYNAMICS

The scalar Lagrangian density $L(\Phi, X)$ is given by

$$L(\Phi, X) = \partial_\mu \Phi \partial^\mu \Phi^* + \partial_\mu X \partial^\mu X^* - V(\Phi, X), \tag{7}$$

where the scalar components of the superfields are denoted by the same symbols as the corresponding superfields. The scalar potential, V, of the chiral superfields $X(x,\theta)$ and $\Phi(x,\theta)$ is given by

$$V = e^K \left\{ \left(\frac{\partial^2 K}{\partial z^i \partial z_j^*} \right)^{-1} D_{z^i} W D_{z_j^*} W^* - 3|W|^2 \right\}, \qquad (z^i = \Phi, X) \tag{8}$$

with

$$D_{z^i} W = \frac{\partial W}{\partial z^i} + \frac{\partial K}{\partial z^i} W. \tag{9}$$

It is explicitly given by

$$V = v^4 e^K \left[\ \left| 1 - g\Phi^2 \right|^2 (1 - |X|^2 + |X|^4) + |X|^2 \left| -2g\Phi + (\Phi + \Phi^*)(1 - g\Phi^2) \ \right|^2 \ \right]. \tag{10}$$

Now, we decompose the scalar field Φ into real and imaginary components,

$$\Phi = \frac{1}{\sqrt{2}} (\varphi + i\chi). \tag{11}$$

Then, the Lagrangian density $L(\varphi, \chi, X)$ is given by

$$L(\varphi, \chi, X) = \frac{1}{2} \partial_\mu \varphi \partial^\mu \varphi + \frac{1}{2} \partial_\mu \chi \partial^\mu \chi + \partial_\mu X \partial^\mu X^* - V(\varphi, \chi, X), \tag{12}$$

with the potential $V(\varphi, \chi, X)$ given by

$$\begin{aligned} V(\varphi, \chi, X) = v^4 \exp(|X|^2 + \varphi^2) \Big[& \left\{ (1 - \frac{g}{2}\varphi^2)^2 + g\chi^2 (1 + \frac{g}{2}\varphi^2 + \frac{g}{4}\chi^2) \right\} (1 - |X|^2 + |X|^4) \\ & + |X|^2 \left\{ 2g^2\chi^2 + 2(g-1)^2\varphi^2 + 2g(g+1)\varphi^2\chi^2 + 2g(g-1)\varphi^4 + \frac{g^2}{2}\varphi^2(\varphi^2 + \chi^2)^2 \right\} \Big]. \end{aligned} \tag{13}$$

While $\varphi, |X| \lesssim \mathcal{O}(1)$ due to the factor $e^{|X|^2 + \varphi^2}$, χ can take a value much larger than unity without costing exponentially large potential energy. Then, the amplitudes of both φ and X soon become smaller than unity due to this steep potential and the exponential factor can be Taylor expanded around the origin. This is the situation we deal with hereafter. Then the scalar potential is dominated by

$$V \simeq \frac{1}{4} \lambda \chi^4, \tag{14}$$

with $\lambda = g^2 v^4$.

Thus chaotic inflation can set out around the Planck epoch. During chaotic inflation, the mass squared of φ, m_φ^2, reads,

$$m_\varphi^2 \simeq \frac{1}{2} g^2 v^4 \chi^4 \simeq 6H^2, \qquad H^2 \simeq \frac{\lambda}{12} \chi^4, \tag{15}$$

where H is the Hubble parameter in this era. Then, φ oscillates rapidly around the origin so that its amplitude damps in proportion to $a^{-3/2}$ with a being the scale factor. At the end of chaotic inflation, the amplitude of φ, $\varphi_{\rm end}$, reads

$$\varphi_{\rm end} \simeq \varphi_{\rm init} e^{-\frac{3}{2}N_{\rm ch}}, \tag{16}$$

classically, where $\varphi_{\rm init}$ is the value of φ in the beginning of chaotic inflation with its natural magnitude being of order of unity, and $N_{\rm ch}$ is the number of e-folds during chaotic inflation which is usually very large. Thus φ practically vanishes classically and its magnitude is dominated by quantum fluctuations. We therefore find $\varphi_{\rm end} \sim H_{\rm end} \sim \lambda^{1/2} = gv^2$.

On the other hand, the mass squared of X, m_X^2, is dominated by,

$$m_X^2 \simeq 2g^2 v^4 \chi^2 \simeq \frac{24}{\chi^2} H^2, \tag{17}$$

which is much smaller than the Hubble parameter in the early stage of chaotic inflation, when X moves towards the origin only slowly. Below we set X to be real and positive making use of the freedom of the phase choice. In this regime classical equations of motion for X and χ are given, respectively, by

$$3H\dot{X} \simeq -m_X^2 X, \tag{18}$$

$$3H\dot{\chi} \simeq -\lambda\chi^3, \tag{19}$$

from which we find,

$$X \propto \sqrt{\chi}. \tag{20}$$

This relation holds actually if and only if quantum fluctuations are unimportant for both χ and X. As for χ, the amplitude of quantum fluctuations acquired in one expansion time is larger than the magnitude of classical evolution in the same period if $\chi \gtrsim \lambda^{-1/6}$ [20], when the universe is in a self-reproduction stage of eternal inflation [20,21]. Hence let us consider the regime $\chi \ll \lambda^{-1/6}$ and (19) holds. Then we can estimate the root-mean-square (RMS) fluctuation in X using the Fokker-Planck equation for the statistical distribution function of X, $P[X,t]$,

$$\frac{\partial}{\partial t}P[X,t] = \frac{1}{3H(t)}\frac{\partial}{\partial X}\left(m_X^2 X P[X,t]\right) + \frac{H^3(t)}{8\pi^2}\frac{\partial^2}{\partial X^2}P[X,t], \tag{21}$$

which is obtained on the basis of (18) using the stochastic inflation method of Starobinsky [22]. Time evolution of the RMS fluctuation of X is given by

$$\frac{d}{dt}\left\langle (\Delta X)^2 \right\rangle = -\frac{2m_X^2}{3H}\left\langle (\Delta X)^2 \right\rangle + \frac{H^3}{4\pi^2}. \tag{22}$$

Taking χ as a time variable in (22) by virtue of (19), we find that the RMS fluctuation of X in an initially homogeneous domain at $\chi = \chi_i$ is given by

$$\left\langle (\Delta X)^2 \right\rangle = \frac{\lambda}{384\pi^2}\left(\chi_i^2 \chi^4 - \chi^6\right), \tag{23}$$

at the epoch χ. Taking $\chi_i \sim \lambda^{-1/6}$, $\left\langle (\Delta X)^2 \right\rangle$ asymptotically approaches

$$\left\langle (\Delta X)^2 \right\rangle \simeq \frac{\lambda^{2/3}}{384\pi^2}\chi^4, \tag{24}$$

which is much less than unity because λ must be a tiny number as will be shown later. From (20) and (24) the amplitude of X becomes much smaller than unity by the time $\chi \simeq \sqrt{24}$. Thereafter (18) no longer holds and X oscillates around the origin rapidly and its amplitude decreases even more. Thus our approximation that both φ and X are much smaller than unity is consistent throughout the chaotic inflation regime.

As χ becomes of order of unity, chaotic inflation ends and the field oscillates coherently with the mass squared $m_\chi^2 \simeq 2gv^4$. Since g must take a value slightly larger than unity as shown later, the energy density of this oscillation becomes less than the vacuum energy density $\sim v^4$ soon. At this stage φ is still localized at the origin since its mass squared is given by

$$m_\varphi^2 \simeq -(g-1) + g\left(1+\frac{g}{2}\right)\chi^2, \tag{25}$$

which becomes negative only after the square amplitude of χ becomes smaller than $\chi_c^2 \equiv 2(g-1)/3 \ll 1$. Thus, the second inflation takes place, which is driven by a false vacuum energy, before φ becomes unstable.

Once the inflation occurs, the amplitude of the oscillation rapidly goes to zero because the amplitude damps in proportional to $a^{-3/2}$. Then, the potential energy with $\chi \simeq 0$ reads,

$$\begin{aligned} V(\varphi) &\simeq v^4 \exp\left(\varphi^2 + |X|^2\right) \\ &\quad\times\left[\ \left(1-\frac{g}{2}\varphi^2\right)^2 (1-|X|^2+|X|^4) + |X|^2\left\{2(g-1)^2\varphi^2 + 2g(g-1)\varphi^4 + \frac{g^2}{2}\varphi^6\right\}\ \right] \\ &\simeq v^4\left[1-\frac{c}{2}\varphi^2 + 2(g-1)^2\varphi^2|X|^2\right] \qquad\qquad (\text{for } \varphi, |X| \ll 1) \\ &\simeq v^4 - \frac{c}{2}v^4\varphi^2, \end{aligned} \tag{26}$$

with $c \equiv 2(g-1)$. Thus, if $g \gtrsim 1$, φ rolls down slowly toward the vacuum expectation value $\eta = \sqrt{2/g}$ and new inflation takes place. Here and hereafter we set φ to be positive by use of Z_2 symmetry($\Phi \leftrightarrow -\Phi$).

As shown in Eq.(16), the initial value of φ in the beginning of new inflation is set by the amplitude of quantum fluctuation and is of order of H then. The dynamics of φ is also governed by quantum fluctuations until the classical motion during the Hubble time, $\Delta\varphi \sim |\dot{\varphi}|H^{-1}$, becomes larger than the quantum fluctuations, $H/(2\pi)$, acquired during the same period. In our case, this condition is equivalent to

$$\varphi > \varphi_c \equiv \frac{v^2}{2\sqrt{3}\pi c} = \frac{H}{2\pi c}. \tag{27}$$

Since $c < 1$ as shown later, quantum fluctuations dominate the dynamics initially. Therefore, the universe enters the self-regenerating stage [20,21] so that the imprint of chaotic inflation is washed away and current horizon scale must be contained in one of a new inflation domain in which φ got larger than φ_c and the classical description of the dynamics of φ with the potential Eq.(26) became possible. The slow-roll condition for the inflaton φ is satisfied for $0 < c < 1$ and $0 \lesssim \varphi \lesssim 1$. The Hubble parameter during new inflation is given by $H \simeq v^2/\sqrt{3}$. Then, the number of e-folds acquired for $\varphi > \varphi_N$ is given by

$$N \simeq \int_{\varphi_f}^{\varphi_N} \frac{V}{V'} \simeq -\frac{1}{c} \ln \left(\frac{\varphi_N}{\varphi_f} \right) \simeq -\frac{1}{c} \ln \varphi_N \tag{28}$$

where the prime represents the derivative with respect to φ and $\varphi_f \sim 1$ is the value of φ at the end of new inflation. In the new inflation regime, both φ and X acquire large quantum fluctuations. Following the same procedure as [23], however, we can show that only φ contributes to adiabatic fluctuations in the present situation. Hence the amplitude of curvature perturbation Φ_H on the comoving horizon scale at $\varphi = \varphi_N$ is given by the standard one-field formula and reads

$$\Phi_H(N) \simeq \frac{f}{2\sqrt{3}\pi} \frac{v^2}{c\varphi_N}, \tag{29}$$

where $f = 3/5$ $(2/3)$ in the matter (radiation) domination. Since the COBE normalization requires $\Phi_H(N) \simeq 3 \times 10^{-5}$ at $N \simeq 60$ [24], the scale v is given by

$$v \simeq 2.3 \times 10^{-2} \sqrt{c} e^{-\frac{cN}{2}}|_{N=60} \simeq 1.8 \times 10^{-3} - 3.6 \times 10^{-4}, \tag{30}$$

for $0.02 \leq c \leq 0.1$. The spectral index n_s of the density fluctuations is given by

$$n_s \simeq 1 - 2c. \tag{31}$$

The data of anisotropies of the cosmic microwave background radiation (CMB) by the COBE satellite [24] also implies $n_s = 1.0 \pm 0.2$ so that $0 < c < 0.1$, which leads to $1.00 < g < 1.05$.

After new inflation, φ oscillates around the minimum $\varphi = \eta$ and the universe is dominated by a coherent scalar-field oscillation of $\sigma \equiv \varphi - \eta$. Since the exponential factor e^{φ^2} in e^K can be expanded as

$$e^{\varphi^2} = e^{\eta^2}(1 + 2\eta\sigma + \cdots),$$

σ has linear interactions of gravitational strength with all scalar and spinor fields in the theory including those in the minimal supersymmetric standard model (MSSM). Let us consider, for example, the term $W = y_i D_i H S_i$ in the superpotential in MSSM, where D_i, S_i are doublet(singlet) superfields, H represents Higgs superfields, and y_i are Yukawa coupling constants. Then, the relevant interaction Lagrangian reads,

$$\mathcal{L}_{\rm int} \sim y_i^2 \eta\sigma D_i^2 S_i^2 + \cdots, \tag{32}$$

which leads to the decay width Γ given by

$$\Gamma \sim \Sigma_i y_i^4 \eta^2 m_\sigma^3. \tag{33}$$

Here $m_\sigma = 2\sqrt{g}e^{1/g}v^2$ is the mass of φ at the vacuum expectation value $\eta = \sqrt{2/g}$. Thus the universe is reheated even if the system of Φ and X has no direct coupling with the standard fields, and the reheat temperature T_R is given by

$$T_R \sim 0.1 \bar{y} \eta m_\varphi^{\frac{3}{2}}, \tag{34}$$

where $\bar{y} = \sqrt{\Sigma_i y_i^4}$. Taking $\bar{y} \sim 1$, the reheating temperature T_R is given by

$$T_R \sim 10^{-10} - 10^{-8}, \tag{35}$$

for $0.02 \leq c \leq 0.1$, which is low enough to avoid the overproduction of gravitinos in a wide range of the gravitino mass [25].

Finally we comment on higher order terms in the Kähler potential. Fourth order terms in the Kähler potential such as

$$\Delta K = \frac{k_1}{2}|X|^2(\Phi + \Phi^*)^2 + \frac{k_2}{4}|X|^4 + \frac{k_3}{12}(\Phi + \Phi^*)^4, \tag{36}$$

roughly generate three minor changes associated with k_i. The first change is as follows; during chaotic inflation, the mass squared m_X^2 of X acquired the additional terms,

$$\begin{aligned} m_X^2 &\simeq 2g^2 v^4 \chi^2 - \frac{k_2}{4} g^2 v^4 \chi^4 \\ &\simeq -3k_2 H^2. \end{aligned} \tag{37}$$

Thus, if $k_2 < -\frac{3}{4}$, X oscillates rapidly so that the amplitude of the oscillation damped into zero. Thus, we can safely set $|X|$ to be zero in the whole analysis of the present model. The second one is the change of $c = 2(g-1)$ into $c = 2(g + k_1 - 1)$. k_3 is almost irrelevant for the dynamics of φ field and only changes its vacuum expectation value η due to the redefinition of φ with a canonical kinetic term.

Furthermore, the inflaton may have interactions with standard light particles ψ_i in the Kähler potential which is invariant under the Nambu-Goldstone-like shift symmetry and $U(1)_R \times Z_2$ symmetry, $(\Phi, \psi_i) = \Sigma_i \frac{\lambda_i}{2}(\Phi + \Phi^*)^2 |\psi_i|^2$. Then, the interaction Lagrangian density is given by $\mathcal{L}_{\rm int} = \Sigma_i \lambda_i \varphi^2 \partial^\mu \psi_i \partial_\mu \psi_i^*$, which yields the decay width Γ of the same order as that in eq.(33) and gives the similar reheating temperature.

IV. CONCLUSIONS AND DISCUSSIONS

In the present paper, we have proposed a consistent scenario of successive inflation in supergravity. Chaotic inflation takes place around the Planck scale so that the universe can live long enough. In this regime, the inflaton responsible for new inflation dynamically relaxes toward zero so that new inflation sets in. The reheating temperature is low enough to avoid overproduction of gravitinos in a wide range of the gravitino mass. Furthermore, our model generally predicts a tilted spectrum with the spectral index $n_s < 1$, which may naturally explain the recent observation of anisotropies of CMB by the BOOMERANG experiment [15] and the MAXIMA experiment [16]. Our model can also accommodate a leptogenesis scenario where the inflaton decay produces heavy Majorana neutrinos and in succession these neutrinos decay to generate the lepton asymmetry to explain the baryon asymmetry observed in our universe [26].

In the present model, the initial value of the inflaton for new inflation may be so close to the local maximum of the potential that the universe enters a self-regenerating stage. Therefore all the scales observable today left the Hubble radius during the last inflation and we cannot verify the chaotic inflation stage directly because the minimum of Kähler potential during chaotic inflation coincides with the local maximum of the potential for new inflation. By appropriately shifting the local minimum of the new inflaton's potential during chaotic inflation one can construct a model in which duration of new inflation is short enough that the trace of chaotic inflation is observable on the large-scale structure. This issue will be discussed in a forthcoming publication.

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